

Radiative corrections in SUSY phenomenology *

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ABSTRACT

We discuss some aspects of the radiative corrections in the phenomenology of the minimal SUSY standard model, by reviewing two recent studies. (1) The full one-loop corrections to the Higgs boson decays into charginos are presented, with emphasis on the renormalization of the chargino sector, including of their mixing matrices. (2) The two-loop $O(\alpha_s \tan \beta)$ corrections to the $b \rightarrow s\gamma$ decay in models with large $\tan \beta$, mainly those to the charged Higgs boson contributions, are discussed. Exact two-loop result is compared to an approximation used in previous studies.

1. Introduction

There are many cases where the radiative corrections become important in the phenomenology of the minimal supersymmetric (SUSY) standard model (MSSM) [1].

(1) Of course, the radiative corrections become large when they are enhanced by large coupling constants and/or large logarithms. For example, QCD corrections to the processes involving quarks, gluon, and their superpartners, are indispensable in the study of the SUSY particles at hadron colliders.

(2) Corrections to the observables which may be precisely measured in present or future experiments are also important. For example, electroweak precision measurements have provided a powerful tool to impose constraints on the SUSY particles. Also, the masses and couplings of several lighter SUSY particles are expected to be precisely measured at future linear colliders [2].

(3) Radiative corrections may generate couplings which are strongly suppressed or even forbidden at lower levels of perturbation. As is well-known, the flavor-changing neutral current (FCNC) is forbidden at the tree-level of the standard model and sensitive to various types of new physics, including the SUSY particles. An example specific to the MSSM is the self-couplings of the Higgs bosons. The SUSY relation between the self-couplings and the electroweak gauge couplings is violated by loop corrections, resulting significant increase of the mass of the lightest Higgs boson h^0 [3] beyond the theoretical upper limit at the tree-level.

In this talk, we review two interesting recent studies of the radiative corrections in the MSSM phenomenology. In section 2, as a case of the class (2) listed above, the full one-loop corrections to the decays of heavier Higgs bosons into charginos are discussed, following Ref. [4]. The role of the renormalization of the chargino sector, including their mixing matrices, is explained in detail. In section 3, as a case of the class (3), two-loop

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$O(\alpha_s \tan \beta)$ SUSY QCD corrections to the $(b \rightarrow s\gamma, b \rightarrow sg)$ decays in models with large $\tan \beta \equiv \langle H_U \rangle / \langle H_D \rangle$, especially the corrections to the charged Higgs boson contribution [5], is discussed. Validity of the approximated calculation of the two-loop integrals, used in previous studies, is examined by comparison with the exact two-loop calculation.

2. One-loop Correction to the Chargino-Higgs boson Couplings

Most of new particles in the MSSM, such as the SUSY particles and Higgs bosons, are mixtures of several gauge eigenstates [1,6]. Mixings of particles therefore play a crucial role in phenomenological studies of these particles.

As an example, the charged SU(2) gauginos \tilde{W}_L^\pm and higgsinos $(\tilde{H}_D^-, \tilde{H}_U^+)_L$ mixes with each other to form two mass eigenstates $\tilde{\chi}_i^\pm (i = 1, 2)$, charginos, as

$$\tilde{\chi}_{iL}^+ = V_{i\alpha} \begin{pmatrix} \tilde{W}_L^+ \\ \tilde{H}_{UL}^+ \end{pmatrix}_\alpha, \quad \tilde{\chi}_{iL}^- = U_{i\alpha} \begin{pmatrix} \tilde{W}_L^- \\ \tilde{H}_{DL}^- \end{pmatrix}_\alpha \quad (i = 1, 2). \quad (1)$$

At the tree-level, the mixing matrices (V, U) are determined to diagonalize the mass matrix

$$X = \begin{pmatrix} M & \sqrt{2}m_W \sin \beta \\ \sqrt{2}m_W \cos \beta & \mu \end{pmatrix} = U^T \begin{pmatrix} m_{\tilde{\chi}_1^+} & 0 \\ 0 & m_{\tilde{\chi}_2^+} \end{pmatrix} V. \quad (2)$$

M and μ are the mass parameters of the SU(2) gaugino and higgsinos, respectively. Couplings of the charginos are generally dependent on (V, U) .

In future colliders, the masses and interactions of the charginos are expected to be measured precisely [2,7,8]. It is therefore very interesting to study the radiative corrections to chargino interactions. In calculating the radiative corrections, we need to renormalize the chargino parameters, including the mixing matrices (V, U) . The renormalization of the chargino sector has been studied for different processes, such as $e^+e^- \rightarrow \tilde{\chi}^+\tilde{\chi}^-$ [9,10,11], $\tilde{f} \rightarrow f'\tilde{\chi}^\pm (f = q, l)$ [12], $H^+ \rightarrow \tilde{\chi}^+\tilde{\chi}^0$ [13], and $\tilde{\chi}^+ \rightarrow \tilde{\chi}^0W^+$ [14].

In this talk, we consider the decays of the heavier Higgs bosons (H^0, A^0) into chargino pair,

$$(H^0, A^0) \rightarrow \tilde{\chi}_i^+ + \tilde{\chi}_j^-, \quad (3)$$

with $i, j = (1, 2)$. If $\tan \beta$ is not much larger than one, the decays (3) may have non-negligible branching ratios [15,16,17]. These decays are also interesting because they are very sensitive to the mixings of charginos. Detailed studies of these decays, including radiative corrections, would therefore provide useful information about the chargino sector, complementary to the pair production processes $e^+e^- \rightarrow \tilde{\chi}_i^+\tilde{\chi}_j^-$ [18,7,8].

The tree-level widths of the decays (3) are $(H_{\{1,2,3\}}^0 \equiv \{h^0, H^0, A^0\})$

$$\begin{aligned} \Gamma^{\text{tree}}(H_k^0 \rightarrow \tilde{\chi}_i^+ \tilde{\chi}_j^-) &= \frac{g^2}{16\pi m_{H_k^0}^3} \kappa(m_{H_k^0}^2, m_{\tilde{\chi}_i^+}^2, m_{\tilde{\chi}_j^+}^2) \\ &\times \left[\left(m_{H_k^0}^2 - m_{\tilde{\chi}_i^+}^2 - m_{\tilde{\chi}_j^+}^2 \right) (F_{ijk}^2 + F_{jik}^2) - 4\eta_k m_{\tilde{\chi}_i^+} m_{\tilde{\chi}_j^+} F_{ijk} F_{jik} \right], \end{aligned} \quad (4)$$

Here $\kappa(x, y, z) \equiv ((x - y - z)^2 - 4yz)^{1/2}$ and η_k is the CP eigenvalue of H_k^0 ($\eta_{1,2} = 1, \eta_3 = -1$). Here we assume that the contributions of CP violation and generation mixings of

the quarks and squarks are negligible. The tree-level couplings gF_{ijk} of the Higgs bosons and charginos $H_k^0 \tilde{\chi}_i^+ \tilde{\chi}_j^+$ come from the gaugino-higgsino-Higgs boson couplings and take the forms [6]

$$gF_{ijk} = \frac{g}{\sqrt{2}} (e_k V_{i1} U_{j2} - d_k V_{i2} U_{j1}) \quad (5)$$

$$\begin{aligned} e_k &= \left(-\sin \alpha, \cos \alpha, -\sin \beta, \cos \beta \right)_k, \\ d_k &= \left(-\cos \alpha, -\sin \alpha, \cos \beta, \sin \beta \right)_k. \end{aligned} \quad (6)$$

Here α is the mixing angle for (h^0, H^0) . The Nambu-Goldstone mode $H_4^0 \equiv G^0$ is included here for later convenience.

The one-loop correction to the coupling gF_{ijk} is expressed as

$$gF_{ijk}^{\text{corr.}} = gF_{ijk} + \delta(gF_{ijk}^{(v)}) + g\delta F_{ijk}^{(w)} + \delta(gF_{ijk}^{(c)}), \quad (7)$$

where $\delta(gF_{ijk}^{(v)})$, $g\delta F_{ijk}^{(w)}$, and $\delta(gF_{ijk}^{(c)})$ are the proper vertex correction, the wave function correction to the external particles, and the counterterm by the renormalization of the parameters (g, V, U, α, β) in the tree-level coupling (6), respectively. The corrections from quarks and squarks in the third generation were calculated in Ref. [19]. Here we present the full one-loop corrections shown in Ref. [4], and show some numerical results for the $(A^0, H^0) \rightarrow \tilde{\chi}_1^+ \tilde{\chi}_1^-$ decays.

We discuss the wave function corrections $\delta F_{ijk}^{(w)}$ in detail. They are expressed as

$$\delta F_{ijk}^{(w)} = \frac{1}{2} \left[\delta Z_{lk}^{H^0} F_{ijl} + \delta Z_{i'i}^{+L} F_{i'jk} + \delta Z_{j'j}^{+R} F_{ij'k} \right]. \quad (8)$$

δZ^{+L} and δZ^{+R} are corrections for the charginos, while $\delta Z_{lk}^{H^0}$ is for the Higgs bosons with $l = (1, 2)$ for $k = (1, 2)$ and $l = (3, 4)$ for $k = 3$. They are given in terms of the self-energies of the relevant particles. Explicit form of $\delta Z_{j'j}^{+L}$, wave function correction to the left-handed chargino $\tilde{\chi}_j^+$, is given by

$$\begin{aligned} \delta Z_{ii}^{+L} &= \\ &- \text{Re} \left\{ \Pi_{ii}^{\tilde{\chi}L}(m_i^2) + m_i \left[m_i \dot{\Pi}_{ii}^{\tilde{\chi}L}(m_i^2) + m_i \dot{\Pi}_{ii}^{\tilde{\chi}R}(m_i^2) + 2 \dot{\Pi}_{ii}^{\tilde{\chi}S,L}(m_i^2) \right] \right\}, \end{aligned} \quad (9)$$

$$\begin{aligned} \delta Z_{pi}^{+L} &= \\ &\frac{2}{m_p^2 - m_i^2} \text{Re} \left\{ m_i^2 \Pi_{pi}^{\tilde{\chi}L}(m_i^2) + m_i m_p \Pi_{pi}^{\tilde{\chi}R}(m_i^2) + m_p \Pi_{pi}^{\tilde{\chi}S,L}(m_i^2) + m_i \Pi_{pi}^{\tilde{\chi}S,R}(m_i^2) \right\}, \end{aligned} \quad (10)$$

where $p \neq i$ and

$$\Pi_{ij}^{\tilde{\chi}}(p) = \Pi_{ij}^{\tilde{\chi}L}(p^2) \not{p} P_L + \Pi_{ij}^{\tilde{\chi}R}(p^2) \not{p} P_R + \Pi_{ij}^{\tilde{\chi}S,L}(p^2) P_L + \Pi_{ij}^{\tilde{\chi}S,R}(p^2) P_R, \quad (11)$$

are the self-energies of the charginos $\tilde{\chi}^+$. δZ^{+R} for the right-handed chargino $\tilde{\chi}_R^+$ is obtained from Eqs. (9, 10) by the exchange $L \leftrightarrow R$. We used the CP symmetry relation

$\text{Re}\Pi_{ii}^{\tilde{\chi}S,L} = \text{Re}\Pi_{ii}^{\tilde{\chi}S,R}$ in Eq. (9). The corrections δZ^{H^0} are

$$\delta Z_{kk}^{H^0} = -\text{Re}\dot{\Pi}_{kk}^{H^0}(m_{H_k^0}^2), \quad k = 1, 2, 3, \quad (12)$$

$$\delta Z_{ab}^{H^0} = \frac{2}{m_{H_a^0}^2 - m_{H_b^0}^2} \text{Re}\Pi_{ab}^{H^0}(m_{H_b^0}^2), \quad a, b = (1, 2), a \neq b \quad (13)$$

$$\delta Z_{43}^{H^0} = -\frac{2}{m_{A^0}^2} \text{Re}\Pi_{43}^{H^0}(m_{A^0}^2). \quad (14)$$

The Higgs boson self-energies $\Pi^{H^0}(k^2)$ in Eqs. (12, 13, 14) include momentum-independent contributions from the tadpole shifts [20] and leading higher-order corrections. The latter contribution is numerically relevant for h^0 and H^0 . Note that $\delta Z_{43}^{H^0}$ in Eq. (14) includes both the $A^0 - G^0$ and $A^0 - Z^0$ mixing contributions.

The off-diagonal part of the wave function correction $\delta Z_{ij}(i \neq j)$ is generated by the mixing between the tree-level mass eigenstates at the one-loop level, and closely related to the renormalization of the mixing matrices. To see this point, we focus on the contribution of δZ_{ii}^{+L} in Eq. (8) and decompose δZ^{+L} into hermitian and anti-hermitian parts, to obtain

$$\begin{aligned} \frac{1}{2}(\delta Z_{ii}^{+L}F_{ijk} + \delta Z_{pi}^{+L}F_{pjk}) &= \frac{1}{2}\delta Z_{ii}^{+L}F_{ijk} + \frac{1}{4}[\delta Z_{pi}^{+L} + (\delta Z_{ip}^{+L})^*]F_{pjk} \\ &\quad + \frac{1}{4}[\delta Z_{pi}^{+L} - (\delta Z_{ip}^{+L})^*]F_{pjk}. \end{aligned} \quad (15)$$

The ultraviolet (UV) divergence of the hermitian part in the first line is cancelled by that of the vertex correction $\delta F_{11k}^{(v)}$ and the counterterm δg . On the other hand, the divergence of the anti-hermitian part in the second line is cancelled by the counterterm δV for the mixing matrix V of $\tilde{\chi}_L^+$, giving

$$\delta F_{ijk}^{(c)}(\delta V) = (\delta V \cdot V^\dagger)_{ip}F_{pjk}. \quad (16)$$

The matrix $\delta V \cdot V^\dagger$ should be anti-hermitian for the unitarity of V and $V^{\text{bare}} \equiv V + \delta V$. Similarly, the UV divergences of the anti-hermitian parts of δZ^{+R} and δZ^{H^0} are cancelled by renormalization of U and α (for H^0, h^0) or β (for A^0), respectively. This relation between the UV divergence of the anti-hermitian part of the wave function corrections δZ and the renormalization of the corresponding mixing matrix holds for general cases [21,22,23].

To fix the chargino sector, we have to specify two input parameters corresponding to two parameters (M, μ) in the mass matrix (2), in addition to $\tan\beta$ which is determined by the Higgs boson sector. The pole masses $m_{\tilde{\chi}_i^+}$ and renormalized mixing matrices $(V, U)^{(\text{ren})}$ are then given as functions of these input parameters. We also need to fix a definition of the renormalized mixing matrices, or the UV finite parts of the counterterms $(\delta V, \delta U)$. In previous studies of the corrections to chargino interactions, several schemes has been proposed for the renormalization of the charginos, as listed below:

(A) We may just use the running mass parameters (M, μ) in the $\overline{\text{DR}}$ scheme at a scale Q as inputs, as in Ref. [9]. Renormalized (V, U) are fixed to diagonalize the tree-level mass matrix. The pole masses $m_{\tilde{\chi}_i^+}$ are shifted from their Q -dependent tree-level values.

The effect of this mass shift has to be taken into account for a proper treatment of the radiative corrections to chargino processes.

(B) On the other hand, one may fix the chargino sector by specifying the pole masses of two charginos $m_{\tilde{\chi}_i^+}$ ($i = 1, 2$), as in Ref. [24]. Renormalized (M, μ) are then defined as tree-level functions of the pole masses. Again, renormalized (V, U) diagonalize the tree-level mass matrix. In this scheme, the pole masses of charginos are identical to their tree-level values by definition. However, one should note that the shift of the masses is unavoidable when the neutralinos $\tilde{\chi}_i^0$ ($i = 1 - 4$) appear in the analysis, since there are only three free parameters (M, μ, M') to describe two charginos and four neutralinos.

(C) Alternatively, we may start from the “on-shell mixing matrices” $(V, U)^{\text{OS}}$ [21,25], defined such that their counterterms completely cancel the anti-hermitian part of the corresponding wave function corrections. For example, the second line of Eq. (15) is dropped by adding Eq. (16) with on-shell δV . The renormalized (M, μ) are then given as diagonal elements of the “on-shell mass matrix” X^{OS} of the charginos [25] given as

$$X^{\text{OS}} = \begin{pmatrix} M^{\text{OS}} & X_{12}^{\text{OS}} \\ X_{21}^{\text{OS}} & \mu^{\text{OS}} \end{pmatrix} \equiv (U^{\text{OS}})^T \begin{pmatrix} m_{\tilde{\chi}_1^+} & 0 \\ 0 & m_{\tilde{\chi}_2^+} \end{pmatrix}_{\text{pole}} V^{\text{OS}} \quad (17)$$

The off-diagonal elements $(X_{12}, X_{21})^{\text{OS}}$ include some information of the loop corrections to the mixings and, as a result, deviate from their tree-level values $\sqrt{2}m_W(\sin \beta, \cos \beta)$. In this scheme, however, both the masses $m_{\tilde{\chi}_i^+}$ and mixing matrices (V, U) are shifted from their tree-level values. Problem from the gauge parameter dependence of the on-shell mixing matrices [26,22,23] may be avoided by improving relevant self energies by the pinch technique [22,27].

We conclude this section with several numerical results, adopted from Ref. [4], for the decay widths of $(A^0, H^0) \rightarrow \tilde{\chi}_1^+ + \tilde{\chi}_1^-$ in the renormalization scheme (C) shown above. Calculation was done by using the packages *FeynArts*, *FormCalc*, and *LoopTools* [28]. We use the SPS1a parameter point [29] as reference point: Chargino and neutralino sectors are specified by the on-shell parameters $M = 197.6$ GeV, $\mu = 353.1$ GeV, $M' = 98$ GeV, and the on-shell parameters for Higgs boson sector, defined as Ref. [20], are $\tan \beta = 10$ and $m_{A^0} = 393.6$ GeV. The SUSY-breaking sfermion-Higgs boson trilinear couplings $(A_t, A_b, A_\tau) = (-487, -766, -250)$ GeV are given in the $\overline{\text{DR}}$ scheme at the parent particle. Other mass parameters for sfermions are $(M_{\tilde{Q}_{1,2}}, M_{\tilde{U}_{1,2}}, M_{\tilde{D}_{1,2}}, M_{\tilde{L}_{1,2}}, M_{\tilde{E}_{1,2}}) = (558.9, 540.5, 538.5, 197.9, 137.8)$ GeV for the first and second generations and $(M_{\tilde{Q}_3}, M_{\tilde{U}_3}, M_{\tilde{D}_3}, M_{\tilde{L}_3}, M_{\tilde{E}_3}) = (512.2, 432.8, 536.5, 196.4, 134.8)$ GeV for the third generation. We used these values in the figures of this section, if not specified otherwise. Using HDECAY program [30], the tree-level branching ratios $\text{Br}(A^0 \rightarrow \tilde{\chi}_1^+ \tilde{\chi}_1^-)$ and $\text{Br}(H^0 \rightarrow \tilde{\chi}_1^+ \tilde{\chi}_1^-)$ at this point are estimated to be 21% and 4%, respectively, which are not negligible.

In Fig. 1, we show the decay widths of $A^0 \rightarrow \tilde{\chi}_1^+ \tilde{\chi}_1^-$ and $H^0 \rightarrow \tilde{\chi}_1^+ \tilde{\chi}_1^-$ as functions of the parent particle, and compare three definitions of the widths: the naive tree-level width Γ^{naive} with the tree-level $m_{\tilde{\chi}_i^+}$ and (V, U) , the tree-level width Γ^{tree} using the pole masses $m_{\tilde{\chi}_i^+}$ and $(V, U)^{\text{OS}}$, and the full one-loop corrected width Γ^{corr} which also includes real photon emission $(A^0, H^0) \rightarrow \tilde{\chi}_1^+ \tilde{\chi}_1^- \gamma$ to cancel infrared divergence. We see that the full one-loop corrections amount up to $\sim -12\%$.

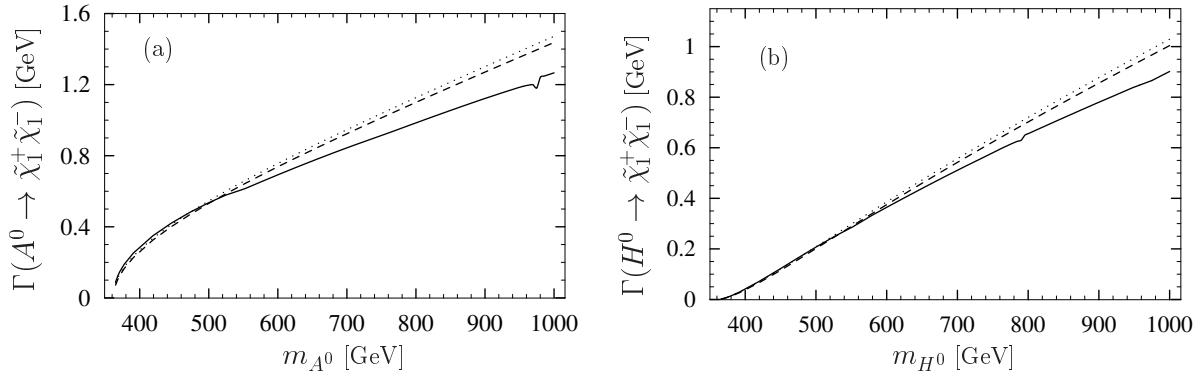


Figure 1: Naive tree-level (dotted), tree-level (dashed), and one-loop corrected (solid) widths of $A^0 \rightarrow \tilde{\chi}_1^+ \tilde{\chi}_1^-$ (a) and $H^0 \rightarrow \tilde{\chi}_1^+ \tilde{\chi}_1^-$ (b) as functions of the parent particle.

In Fig. 2 we compare the contributions from the (s)fermion loops [19] (loops with quarks, leptons, and their superpartners) and the full one-loop contributions, relative to Γ^{naive} , for Fig. 1(a). Corrections to the chargino mass matrix are shown by the dash-dotted line for the (s)fermion loops while the dotted line is for the full correction. The solid (dashed) line shows the total correction including full ((s)fermion) one-loop contributions. Figure 2 shows that the (s)fermion loop corrections and other corrections are of comparable order, both for the chargino mass matrix and for the conventional loop corrections (7).

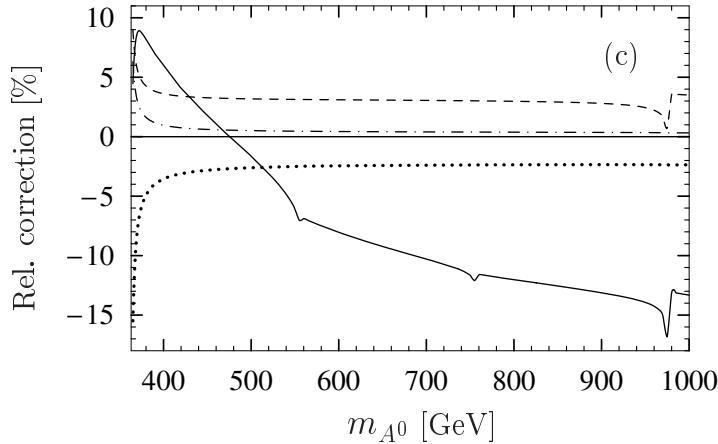


Figure 2: Individual one-loop corrections to the decay width of $A^0 \rightarrow \tilde{\chi}_1^+ \tilde{\chi}_1^-$ relative to the naive tree-level width. Explanation of each line is seen in the text.

Fig. 3 shows the corrections to the decay widths of $(A^0, H^0) \rightarrow \tilde{\chi}_1^+ + \tilde{\chi}_1^-$ as a function of $A_t = A_b = A_\tau$, with the other parameters unchanged. The dashed lines denote $\Gamma^{\text{tree}}/\Gamma^{\text{naive}} - 1$ and show the effect of the chargino mass matrix correction. The solid lines show the total correction $\Gamma^{\text{corr}}/\Gamma^{\text{naive}} - 1$. The dotted lines stand for $\Gamma^{\text{corr}}/\Gamma^{\text{tree}} - 1$, the conventional loop correction in Eq. (7). One sees that the A_t dependence of the cor-

rected widths mainly comes through the shifts of the masses and mixing matrices of the charginos from the tree-level values.

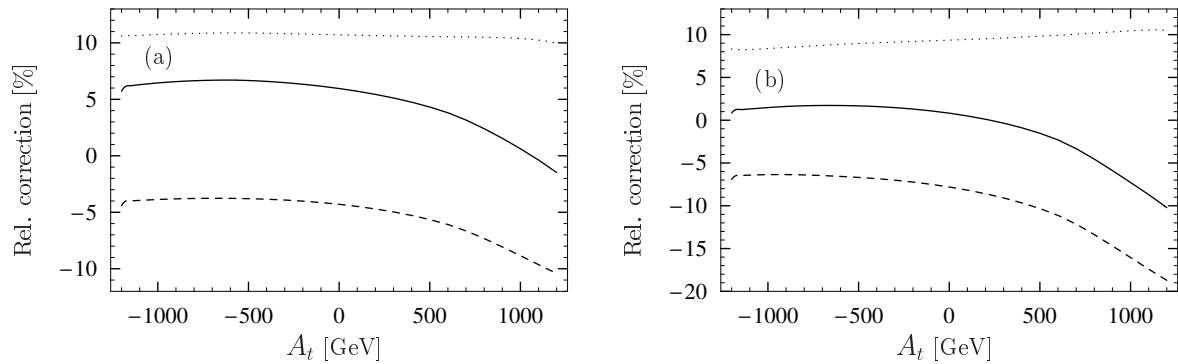


Figure 3: Relative corrections for the decays $A^0 \rightarrow \tilde{\chi}_1^+ + \tilde{\chi}_1^-$ (a) and $H^0 \rightarrow \tilde{\chi}_1^+ + \tilde{\chi}_1^-$ (b) as functions of $A_t = A_b = A_\tau$. The dashed lines, solid lines, and dotted lines denote $\Gamma^{\text{tree}}/\Gamma^{\text{naive}} - 1$, $\Gamma^{\text{corr}}/\Gamma^{\text{naive}} - 1$, and $\Gamma^{\text{corr}}/\Gamma^{\text{tree}} - 1$, respectively.

3. Two-loop $O(\alpha_s \tan \beta)$ Corrections to $b \rightarrow s\gamma$

There are many cases where two-loop and even higher order radiative corrections are necessary in the MSSM phenomenology. As a well-known example, the correction to the mass of the lightest Higgs boson h^0 is so large that the two-loop contribution [31,32] is still larger than the expected error in future measurements.

Here we consider the two-loop $O(\alpha_s \tan \beta)$ SUSY QCD corrections to the $b \rightarrow s\gamma$ and $b \rightarrow sg$ decays in models with large $\tan \beta$. These decays describe the inclusive decay width $\text{Br}(\bar{B} \rightarrow X_s \gamma)$ very well [33], up to the nonperturbative hadronic corrections which are small and well under control.

In the standard model, the decays $b \rightarrow (s\gamma, sg)$ occur through W^\pm boson loops. These decays are important to prove possible new physics beyond the standard model since the new physics may contribute at the same level of perturbation as the standard model one.

In the MSSM, these decays receive new contributions [34,35] from loops with the charged Higgs boson H^\pm , charginos $\tilde{\chi}^\pm$, gluino \tilde{g} , and neutralino. Their contributions are often comparable to or even larger than the W^\pm loop, and sensitive to the masses and couplings of these new particles. The leading order QCD corrections to these new contributions have been calculated [36] for generic models. Higher-order QCD and SUSY QCD corrections have been evaluated for specific models [37,38,39].

Here we are interested in the $b \rightarrow (s\gamma, sg)$ decays in models with very large $\tan \beta$ [40,41]. One important finding is that the SUSY QCD may induce $\mathcal{O}(\alpha_s \tan \beta)$ corrections [38,39] to the contributions of the charged Higgs boson and of charginos. These two-loop corrections may be comparable to the leading one-loop contributions, as shown below, and significantly affect the experimental constraints [40,41] on the new particles. In this talk, we mainly consider the corrections to the contribution of the charged Higgs boson H^+ , following Ref. [5].

At the one-loop level, dominant contribution by H^\pm exchange comes from the diagram in Fig. 4 with initial b_R . Couplings of H^\pm to quarks are derived from the tree-level

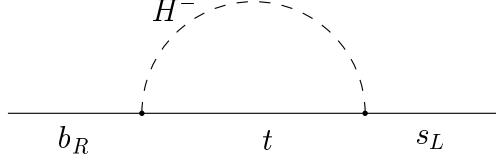


Figure 4: $b \rightarrow s\gamma$ and $b \rightarrow sg$ decays by the one-loop H^\pm exchange. The photon or gluon is to be attached at the t or H^- lines.

lagrangian

$$\mathcal{L}_{\text{int}} = -h_b \bar{b}_R q_L H_D - h_t \bar{t}_R q_L H_U + (\text{h.c.}), \quad (18)$$

where only quarks in the third generation (t, b) are included for simplicity. At the tree-level, the couplings of H_U to b_R and of H_D to t_R are forbidden by SUSY and the Peccei-Quinn symmetry under $(q_L, t_R, H_U) \rightarrow (q_L, t_R, H_U)$, $(b_R, H_D) \rightarrow (-b_R, -H_D)$. However, squark-gluino loops with breakings of both symmetries may induce effective couplings [42,43,44]

$$\Delta \mathcal{L}_{\text{eff.int}} = -h_b \Delta_b \bar{b}_R q_L H_U - h_t \Delta_t \bar{t}_R q_L H_D + (\text{h.c.}). \quad (19)$$

Δ_q ($q = b, t$) are one-loop functions of $O(\alpha_s \mu m_{\tilde{g}}/M_{\text{SUSY}}^2)$, where squarks and gluino masses are around the scale M_{SUSY} . There are also $O(h_t^2)$ contributions to Eq. (19) from squark-higgsino loops. Note that Δ_q do not decouple in large M_{SUSY} limit [42].

Although $|\Delta_q|$ themselves are sufficiently smaller than unity, their contributions to the H^+ couplings are enhanced by $\tan \beta$ relative to the tree-level, as shown below, and may give large corrections in large $\tan \beta$ models.

(i) Correction from counterterm to m_b [42]: The QCD running mass $m_b(\text{SM})$ within the standard model is given by Eqs. (18, 19) as

$$m_b(\text{SM}) = \frac{h_b \bar{v}}{\sqrt{2}} \cos \beta [1 + \Delta_b \tan \beta] \quad (20)$$

$$= m_b(\text{MSSM}) + \delta m_b. \quad (21)$$

The squark-gluino correction δm_b lift tree-level suppression of m_b by $\cos \beta$ and may become comparable to the tree-level contribution. As a result, the $H^+ \bar{t}_L b_R$ coupling y_b may significantly deviate from the tree-level as

$$\begin{aligned} y_b(H^+ \bar{t}_L b_R)(\text{eff}) &= V_{tb} h_b \sin \beta (1 - \Delta_b \cot \beta) \\ &\rightarrow V_{tb} \frac{\sqrt{2} m_b(\text{SM})}{\bar{v}} \tan \beta \frac{1}{1 + \Delta_b \tan \beta}. \end{aligned} \quad (22)$$

The large correction $\Delta_b \tan \beta$ is originated from that the $H^+ \bar{t}_L b_R$ coupling receives very small contribution from Eq. (19) because of $H^+ = \sin \beta H_D^+ + \cos \beta H_U^+ \sim H_D^+$. Similarly, δm_b in Eq. (21) also induce $O(\alpha_s \tan \beta)$ corrections to the couplings of b_R to heavier Higgs bosons (H^0, A^0) and to the higgsino \tilde{H}_D .

(ii) Correction to the $H^- \bar{b}_L t_R$ coupling y_t comes from Δ_t through the proper vertex correction as [43,44,45]

$$y_t(H^+ \bar{b}_L t_R)(\text{eff}) = V_{tb}^* h_t \cos \beta (1 - \Delta_t \tan \beta) \quad (23)$$

$$\rightarrow V_{tb}^* \frac{\sqrt{2} m_t}{\bar{v}} \cot \beta (1 - \Delta_t \tan \beta). \quad (24)$$

In general, Eq. (19) has mixing terms between quarks in different generations, which are induced by the squark-higgsino loops and squark-gluino loops with squark generation mixings [42,45]. These mixing terms may generate $\tan \beta$ -enhanced corrections to the CKM matrix V and flavor-changing couplings of (H^0, A^0) . The latter couplings induce the decays $B_s \rightarrow \mu^+ \mu^-$ [46,45] and $(H^0, A^0) \rightarrow b\bar{s}$ [47].

Two-loop $\mathcal{O}(\alpha_s \tan \beta)$ corrections to the $b \rightarrow (s\gamma, sg)$ decays has been calculated in Refs. [37,38,39,45]. Here we discuss the H^\pm contributions to the Wilson coefficients $C_i(\mu)$ ($i = 7, 8$), defined in the effective Hamiltonian

$$H_{\text{eff}} \supset -\frac{4G_F}{\sqrt{2}} V_{ts}^* V_{tb} (C_7(\mu) \mathcal{O}_7(\mu) + C_8(\mu) \mathcal{O}_8(\mu)), \quad (25)$$

with

$$\mathcal{O}_7(\mu) = \frac{e}{16\pi^2} m_b(\mu) \bar{s}_L \sigma^{\mu\nu} b_R F_{\mu\nu}, \quad \mathcal{O}_8(\mu) = \frac{g_s}{16\pi^2} m_b(\mu) \bar{s}_L \sigma^{\mu\nu} T^a b_R G_{\mu\nu}^a. \quad (26)$$

The H^\pm contributions $C_{i,H}$ ($i = 7, 8$) to $\mathcal{O}(\alpha_s \tan \beta)$ at the scale $\mu_W = m_W$ are expressed as

$$C_{i,H}(\mu_W) = \frac{1}{1 + \Delta_{b_R,b} \tan \beta} [C_{i,H}^0(\mu_W) + \Delta C_{i,H}^1(\mu_W)]. \quad (27)$$

Here $C_{i,H}^0(\mu_W)$ and $\Delta C_{i,H}^1(\mu_W)$ are the contributions of the one-loop diagram and the two-loop diagrams in Fig. 5, respectively. The overall factor $1/(1 + \Delta_{b_R,b} \tan \beta)$ represents the correction from δm_b . The one-loop integral $\Delta_{b_R,b}$ improves Δ_b in Eq. (19) by inclusion of the $SU(2) \times U(1)$ breaking for the masses and couplings of squarks [38,45,5].

In previous studies [37,38,39], the $\mathcal{O}(\alpha_s \tan \beta)$ SUSY QCD corrections were evaluated in terms of an effective two-Higgs-doublet lagrangian, in which squarks and gluino are integrated out. This approach is called the “nondecoupling approximation” in Ref. [5] since it preserves all $\mathcal{O}(M_{\text{SUSY}}^0)$ contributions of the original two-loop integrals. For the corrections (i) from δm_b , this approximation allows us to resum higher-order $\mathcal{O}((\alpha_s \tan \beta)^n)$ terms (Carena et al. in Ref. [43]), by putting $\Delta_{b_R,b}$ in the denominator as in Eq. (27). In contrast, for the proper vertex corrections (ii) to the $H^- \bar{s}_L t$ coupling in Fig. 5, the nondecoupling approximation retains only diagrams (a) and (b), and the squark-gluino subloops are evaluated at vanishing external momenta. The $\mathcal{O}(\alpha_s \tan \beta)$ result (27) is then approximated by a rather simple form

$$C_{i,H}(\mu_W)|_{\text{nondec}} = \frac{1 - \Delta_{t_R,s} \tan \beta}{1 + \Delta_{b_R,b} \tan \beta} C_{i,H}^0(\mu_W). \quad (28)$$

The one-loop integral $\Delta_{t_R,s}$, defined in Ref. [5], corresponds to Δ_t in Eq. (19) applied for the $H^- \bar{s}_L t_R$ coupling, including the $SU(2) \times U(1)$ breaking effect [45].

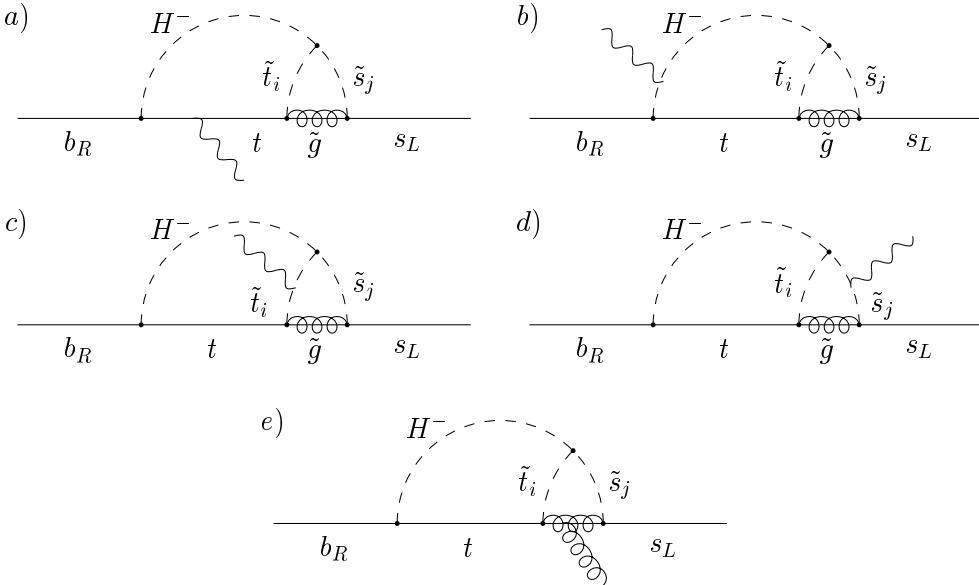


Figure 5: H^\pm mediated diagrams contributing at order $\mathcal{O}(\alpha_s \tan \beta)$ to the decays $b \rightarrow s\gamma$ and $b \rightarrow sg$. The photon must be replaced by a gluon and vice versa, whenever possible.

However, the momentum dependence of the squark-gluino subloops in Fig. 5a,b), as well as the diagrams in Fig. 5c-e) ignored in the nondecoupling approximation, are expected to give $\mathcal{O}((m_{\text{weak}}^2, m_{H^\pm}^2)/M_{\text{SUSY}}^2)$ contributions, where $m_{\text{weak}} \sim (m_W, m_t)$, and, therefore, cause significant deviation of the exact two-loop result from the nondecoupling approximation when M_{SUSY} is not much larger than m_{weak} and/or m_{H^\pm} . It is important to examine, in such cases, how large the deviation is and how far the nondecoupling approximation may be applied beyond the restriction $(m_{\text{weak}}^2, m_{H^\pm}^2) \ll M_{\text{SUSY}}^2$.

We perform an exact evaluation of the two-loop diagrams in Fig. 5 and compare the results to those in the nondecoupling approximation. In Fig. 6, we show the numerical results of $C_{i,H}(\mu_W)$ as functions of m_{H^\pm} , for a SUSY particle spectrum $(m_{\tilde{s}_L}, M_{\tilde{Q}_3}, M_{\tilde{U}_3}, M_{\tilde{D}_3}) = (250, 230, 210, 260)$ GeV, $A_t = 70$ GeV, $A_b = 0$, $\tan \beta = 30$, $m_{\tilde{g}} = 200$ GeV, and $\mu = 250$ GeV. We see that the $\mathcal{O}(\alpha_s \tan \beta)$ corrections are numerically comparable to the one-loop results and must be included in realistic analysis. The deviation of the exact two-loop results is $\mathcal{O}(m_{\text{weak}}^2/M_{\text{SUSY}}^2)$, the same order as the $\text{SU}(2) \times \text{U}(1)$ breaking effects in the squark-gluino subloops [45], and not negligible, especially for $C_{8,H}$. However, contrary to the naive expectation, the deviation does not show significant increase for $m_{H^\pm} > M_{\text{SUSY}}$. This is more clearly seen in the left plot of Fig. 7 where the relative difference between the exact two-loop result and the nondecoupling approximation,

$$r_i(\mu_W) \equiv \frac{C_{i,H}(\mu_W)|_{\text{nondec}} - C_{i,H}(\mu_W)|_{\text{exact}}}{C_{i,H}(\mu_W)|_{\text{exact}}} \quad (i = 7, 8), \quad (29)$$

is shown. For reference, the right plot of Fig. 7 shows the results for a heavier SUSY spectrum $(m_{\tilde{s}_L}, M_{\tilde{Q}_3}, M_{\tilde{U}_3}, M_{\tilde{D}_3}) = (700, 450, 435, 470)$ GeV, $A_t = 150$ GeV, $A_b = 0$, $\tan \beta = 30$, $m_{\tilde{g}} = 600$ GeV, and $\mu = 550$ GeV. r_i is very small in the whole range of m_{H^\pm} . In both cases, the main part of the deviation comes from the diagram in Fig. 5a)

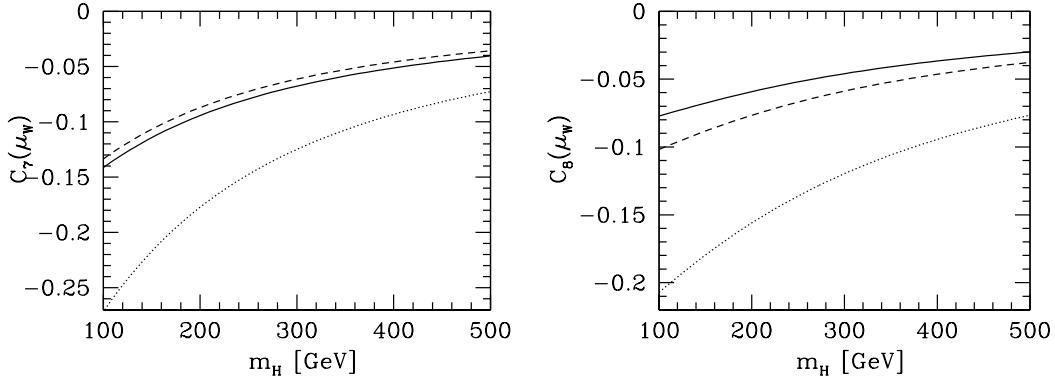


Figure 6: $C_{7,H}(\mu_W)$ and $C_{8,H}(\mu_W)$ as functions of m_H . The dotted, dashed, and solid lines show the one-loop result, nondecoupling approximation, and exact two-loop result, respectively. Parameters for the SUSY particles are shown in the text.

and, for $C_{8,H}$, also from the diagram in Fig. 5e).

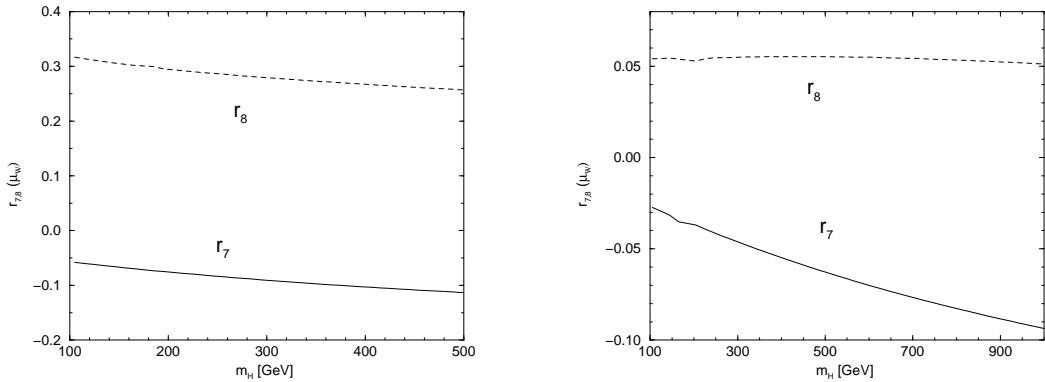


Figure 7: Relative difference $r_i(\mu_W)$ ($i = 7, 8$) between the exact two-loop results and the nondecoupling approximations of $C_{i,H}(\mu_W)$, for the SUSY spectrum as in Fig. 6 (left) and heavier spectrum (right).

To understand this unexpected result for $m_{H^\pm} > M_{\text{SUSY}}$ qualitatively, we consider the diagram (a) in Fig. 5, with chirality flip on the top quark line. When m_{H^\pm} is sufficiently larger than m_t , this diagram gives the largest contribution to $\Delta C_{i,H}^1(\mu_W)$. The contribution is proportional to the loop integral

$$\mu m_{\tilde{g}} I_{ti2}(m_t, m_{H^\pm}, m_{\tilde{t}_i}, m_{\tilde{s}}, m_{\tilde{g}}) = \int \frac{d^4 k}{(2\pi)^4} \frac{k^2}{[k^2 - m_t^2]^3 [k^2 - m_{H^\pm}^2]} Y_{ti2}(k^2; m_{\tilde{t}_i}, m_{\tilde{s}}, m_{\tilde{g}}), \quad (30)$$

where $Y_{ti2}(k^2; m_{\tilde{t}_i}, m_{\tilde{s}}, m_{\tilde{g}})$ represents the squark-gluino subdiagram contribution to the

effective vertex $H^- \bar{s}_L t_R$ and is given by

$$Y_{ti2}(k^2; m_{\tilde{t}_i}, m_{\tilde{s}}, m_{\tilde{g}}) = \mu m_{\tilde{g}} \left[-2F + (k^2 - m_t^2)G \right] (k^2; m_{\tilde{t}_i}^2, m_{\tilde{s}}^2, m_{\tilde{g}}^2), \quad (31)$$

with

$$F(k^2; m_{\tilde{t}_i}^2, m_{\tilde{s}}^2, m_{\tilde{g}}^2) = \int \frac{d^4 l}{(2\pi)^4} \frac{1}{[(l+k)^2 - m_{\tilde{t}_i}^2][l^2 - m_{\tilde{s}}^2][l^2 - m_{\tilde{g}}^2]}, \quad (32)$$

$$k^\mu G(k^2; m_{\tilde{t}_i}^2, m_{\tilde{s}}^2, m_{\tilde{g}}^2) = \int \frac{d^4 l}{(2\pi)^4} \frac{l^\mu}{[(l+k)^2 - m_{\tilde{t}_i}^2][l^2 - m_{\tilde{s}}^2][l^2 - m_{\tilde{g}}^2]^2}. \quad (33)$$

In the nondecoupling approximation, the form factor $Y_{ti2}(k^2; m_{\tilde{t}_i}, m_{\tilde{s}}, m_{\tilde{g}})$ in Eq. (30) is replaced by

$$Y_{ti2}|_{\text{nondec}} = -2\mu m_{\tilde{g}} F(0; m_{\tilde{t}_i}^2, m_{\tilde{s}}^2, m_{\tilde{g}}^2), \quad (34)$$

which is independent of k^2 . For simplicity, we hereafter set $m_{\tilde{t}_i}$, $m_{\tilde{s}}$, $m_{\tilde{g}}$, and μ equal to M_{SUSY} .

For $|k^2|$ much smaller or larger than M_{SUSY}^2 , $Y_{ti2}(k^2; M_{\text{SUSY}}^2)$ behaves as

$$Y_{ti2}(k^2; M_{\text{SUSY}}^2) \rightarrow \begin{cases} Y_{ti2}|_{\text{nondec}} + O\left(\frac{k^2}{M_{\text{SUSY}}^2}, \frac{m_t^2}{M_{\text{SUSY}}^2}\right) & (|k^2| \ll M_{\text{SUSY}}^2), \\ O\left(\frac{M_{\text{SUSY}}^2}{k^2} \ln \frac{k^2}{M_{\text{SUSY}}^2}\right) & (|k^2| \gg M_{\text{SUSY}}^2), \end{cases} \quad (35)$$

which supports the naive expectation that a substantial deviation of $I_{ti2}(m_t, m_{H^\pm}, M_{\text{SUSY}}^2)$ from $I_{ti2}(m_t, m_{H^\pm}, M_{\text{SUSY}}^2)|_{\text{nondec}}$ may arise from the region $|k^2| > M_{\text{SUSY}}^2$.

However, the factor multiplying $Y_{ti2}(k^2; M_{\text{SUSY}}^2)$ in Eqs. (30) drops as $d^4 k / k^6$ for $|k^2| \gg m_{H^\pm}^2$. In fact, the bulk of the integral I_{ti2} is determined by the small $|k^2|$ region up to $|k^2| = O(m_t^2)$. If M_{SUSY} is sufficiently larger than m_t , $Y_{ti2}(k^2; M_{\text{SUSY}}^2)$ does not deviate substantially from $Y_{ti2}|_{\text{nondec}}$ in this region. This explains the smallness of the deviation for $m_{H^\pm} > M_{\text{SUSY}}$ shown in Figs. 6 and 7.

We comment on the SUSY QCD corrections to other one-loop contributions to the $b \rightarrow (s\gamma, sg)$ decays. As already mentioned, the chargino contributions receive the $O(\alpha_s \tan \beta)$ correction to the $\tilde{\chi}^+ b_R \tilde{t}_L^*$ coupling [38,39] from δm_b in Eq. (21), through the coupling h_b of the higgsino component \tilde{H}_D of $\tilde{\chi}^+$ to b_R . Other gluino corrections are not enhanced by $\tan \beta$ relative to the one-loop contribution. In contrast, the W^\pm contributions do not receive $O(\alpha_s \tan \beta)$ corrections in the nondecoupling approximation. However, two-loop diagrams with effective $W^+ \bar{b}_R$ or $G^+ \bar{b}_R$ couplings, some of which are shown in Fig. 8, give decoupling $O(\alpha_s \tan \beta m_{\text{weak}}^2 / M_{\text{SUSY}}^2)$ contributions and may become nonnegligible for light $M_{\text{SUSY}} \sim m_{\text{weak}}$. Numerical study of these contributions will be presented in Ref. [48].

In addition, there are also contributions to $b \rightarrow s\gamma$ coming from the mixings of squarks \tilde{b} and \tilde{s} , such as the one-loop squark-gluino contribution [35,40,36]. One should note that squark generation mixings may be induced by the running of squark mass parameters [49], $O(\tan \beta)$ corrections to the quark Yukawa matrices [50], corrections to the squark- (γ, g)

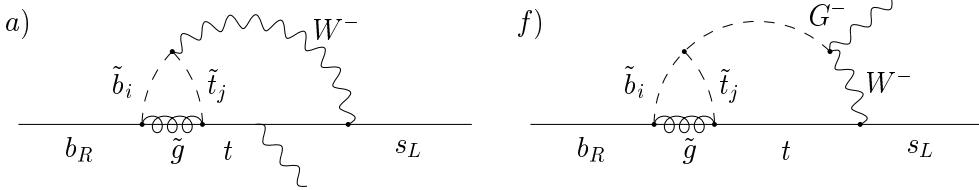


Figure 8: Examples of the two-loop diagrams for the $O(\alpha_s \tan \beta)$ corrections to the W^\pm contributions to $b \rightarrow s\gamma$.

couplings [51], and other loop corrections. Studies of such contributions need consistent treatment of the squark sector renormalization including generation mixings, similar to the discussion in Section 2.

4. Conclusion

We have discussed some aspects of the radiative corrections in the MSSM phenomenology, using two recent studies. First, the full one-loop corrections to the Higgs boson decays into charginos were presented. Especially, the renormalization of the chargino sector, including their mixing matrices, was discussed in detail. Numerical result was shown for the $(A^0, H^0) \rightarrow \tilde{\chi}_1^+ \tilde{\chi}_1^-$ decays. Second, the two-loop $O(\alpha_s \tan \beta)$ corrections to the $b \rightarrow s\gamma$ and $b \rightarrow sg$ decays were discussed in models with large $\tan \beta$. Validity of the nondecoupling approximation, used in previous calculations, was examined for the H^\pm contribution, by exact evaluation of the two-loop diagrams. The deviation was shown to be $O(m_{\text{weak}}^2/M_{\text{SUSY}}^2)$, but, contrary to naive expectation, not increase as m_{H^\pm} even for $m_{H^\pm} > M_{\text{SUSY}}$. A qualitative explanation for this unexpected behavior was presented in terms of the structure of the relevant two-loop integral.

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